Large D-terms, hierarchical soft spectra and moduli stabilisation

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We derive general expressions for soft terms in supergravity where D-terms contribute significantly to the supersymmetry breaking in addition to the standard F-type breaking terms. Such D-terms can strongly influence the scalar mass squared terms, while having limited impact on gaugino masses and the B-terms. We present parameterisations for the soft terms when D-terms dominate over F-terms or become comparable with them. Novel patterns emerge which can be tested phenomenologically. In a mixed anomaly-D mediated scenario, the scalars have masses from D-mediation, whereas gaugino masses are generated by anomaly mediation. As an application of this analysis, we show that while the "split supersymmetry" like mass spectrum with one fine tuned Higgs is not an automatic outcome of these scenarios, explicit models can be constructed where it can be realised. Finally, we show that large D-mediated supersymmetry breaking can be realised in string models based on intersecting D-branes. Examples are presented where the moduli are stabilised in the presence of large D-terms using non-perturbative gaugino condensation like effects.

I. INTRODUCTION

In most models of supersymmetry (SUSY) breaking, supersymmetry is broken spontaneously in a hidden sector and is then transmitted to the visible sector through some interactions, mostly gravitational. In supergravity, the hidden sector typically contains a set of chiral fields whose auxiliary components attain a *vev* at the minimum breaking supersymmetry spontaneously. It is generally preferred to have a dynamical explanation to this phenomenon. This breaking is communicated to the visible sector through tree level (and higher order) gravitational interactions. After integrating out the heavy fields, including the hidden sector fields, the resulting effective lagrangian contains renormalisable supersymmetry breaking soft terms [1, 2]. At the full supergravity (SUGRA) level,

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the soft terms are typically given in terms of the gauge kinetic function f, the Kähler potential K and the superpotential W. Thus in the global limit, the structure of the soft terms crucially depends on the forms these functions take in SUGRA. For example, if the Kähler and the gauge kinetic functions are canonical, this will lead to a universal soft spectrum with mSUGRA boundary conditions.

While analysis of the above type are suitable for simplest classes of supersymmetry breaking models, for more complex situations it is useful to have general expressions for soft terms [3, 4]. Such situations can typically arise when supersymmetry breaking has its origins in string theory. Given that we do not yet have a concrete model of supersymmetry breaking in string theory, it is much more advantageous to parameterise this breaking in terms of a few parameters. In terms of effective supergravity lagrangians derived from string theory, the breaking can be parameterised as the *vev*s of the auxiliary fields of the chiral superfields associated with the higher dimensional gravitational multiplet, namely the dilaton field S and the moduli fields T_i, which effectively act as hidden sector fields. The main advantage of such parameterisations is that they could capture the generic features of soft spectrum emanating from a class of models without completely resorting to explicit model building. These features could then be contrasted with the phenomenological requirements. Detailed analysis parametrising the resultant soft terms for the heterotic case have been presented in [5]. Recently, they have been further extended to the case of Type-I strings [6].

The above analysis which has been very useful can however, be considered as incomplete. This is because, they have implicitly assumed only F-type breaking of supersymmetry (only auxiliary fields of the chiral multiplets get a *vev*). In a more generic scenario, it is well known that there could be D-type susy breaking contributions too [7]. These can arise for example in models based on anomalous $U(1)$ symmetries [8]. Furthermore, in effective lagrangians from the Type II orientifolds with intersecting D-branes, one can expect such D-term contributions to be naturally present. Given these motivations, it is natural to extend the previous analysis by considering D-type SUSY breaking sources. In section two, we present these general expressions of soft terms, initially for generic fields, then for the specific case of the matter fields.

That D-type could have strong impact on the pattern of soft masses has been known for some time, particularly for the limit when the D-terms are small (less than the corresponding F-type contributions), $\lesssim \mathcal{O}(m_{3/2}^2)$ [7]. A more dramatic impact could be expected if the D-terms are large and as allowed by the cosmological constant limit, within the range, $\mathcal{O}(m_{3/2}^2) \lesssim D \lesssim \mathcal{O}(m_{3/2}M_{Pl})$. For example, considering only pure F-type breaking, leads to a typical spectrum of the soft masses, where the gaugino and the Higgsino masses are roughly proportional to the gravitino mass, $m_{3/2}$,

whereas the scalar mass squared and the B-terms are proportional to $m_{3/2}^2$. Adding large D-type sources could significantly alter this simple pattern by generating a splitting between the fermionic and scalar superpartners, by an amount proportional to D. In the extreme limit, this would mean that the scalars can have masses close to the intermediate scale. Following the works of Ref.[5, 6], we parameterise the soft terms in three particular cases (section III) : (i) mixed D and anomaly mediation (ii) mixed D and S mediation and (iii) mixed D and T mediation. In the mixed D and anomaly mediated scenario, scalar masses can be everywhere between the weak scale and an intermediate scale whereas the gaugino masses, B-term and the μ term are proportional to the gravitino mass, $m_{3/2}$, which can be taken close to the weak scale. In the mixed D and S(T) mediated scenarios, the hierarchy between scalar and fermionic superpartners is parameterised by an angle $\gamma_{S(T)}$, which could be constrained by phenomenology.

The splitting due to the D-terms could well have another important application in understanding the origins of recently proposed "split supersymmetry" models. Influenced by multivacua structure in string theory as a possible new view on the cosmological constant problem [9], these models question the solution of the gauge hierarchy problem through low energy SUSY [10]. In this proposal, not all superpartners are required to be at a scale close to TeV. Instead, it is sufficient if the fermionic superpartners stay close to the weak scale, whereas the scalar superpartners can be present at scales as high as 10^9 GeV. This way, one keeps the nice features of gauge coupling unification and the viable dark matter candidate of low energy supersymmetry, while getting rid of unwanted features associated with large flavour changing neutral current effects and CP violation problems[11]. In section IV, we address the question of attaining split supersymmetry by including D-mediation. As we will see, though it is not automatic to have exact split spectrum in these models, specifically due to the B term, we can nevertheless envisage models where it is possible to generate hierarchical spectrum and we will present explicit models of this type.

So far we have not addressed the issue of the origin of such large D-terms. We address this issue in sections V and VI. Unlike in the heterotic case, in Type I/II string theories, Fayet-Iliopoulos terms can appear at the tree level and thus it is possible to generate SUSY breaking with large D-terms. We will present an explicit example in the context of intersecting D-branes Type I orientifold models with four stacks of D9 branes, each stack containing four coincident branes. However, a related question concerns the stabilisation of the moduli as these FI terms are field-dependent. We find that standard mechanisms like gaugino condensation, suitably combined with other mechanisms of moduli stabilisation as, e.g. three-form fluxes in IIB orientifolds, are still applicable even in the limit of large D-terms. We present an example detailing this point. We close

with a summary. A preliminary version of our results was reported in [12].

II.   GENERAL EXPRESSIONS INCLUDING D-BREAKING

In the following, we will present general expressions for the soft terms including D-type supersymmetry breaking terms. As is the case with any general analysis, we will not address the question of the origins of these SUSY breaking *vevs* for either F-terms or D-terms. We will assume SUSY to be broken with both these types of breaking and proceed to derive the soft terms. As a starting point, let us recall the form the scalar potential in supergravity[1]:

$$V = e^G(G^M \ G_M - 3) + \frac{1}{2}\sum_A g_A^2 D_A^2. \tag{1}$$

Here $G = K + \ln|W|^2$, with K being the Kähler potential and W, the superpotential and $1/g_A^2 = Ref_A$, where f_A is the gauge kinetic function. The F terms in the scalar potential are given by $G_M = \partial G/\partial z^M$, where z represents the scalar part of a chiral superfield. The index M runs over all the chiral superfields present, matter as well as hidden sector and/or moduli fields. The D-terms, D_A carry the obvious notation with the index A running over all the $U(1)$ factors present[2].

While deriving the soft terms, a couple of constraints need to be satisfied. First, at the minimum, both D and F terms contribute to supersymmetry breaking and thus to the vacuum energy. This can be canceled by the superpotential (W) *vev* which gives mass to the gravitino. We will impose this fine-tuning condition on the potential. This means:

$$< V > = < e^G(G^M \ G_M - 3) + \frac{1}{2}\sum_A g_A^2 D_A^2 > = \ 0 \ . \tag{2}$$

Second is the necessary condition for the existence of the minima: $< \partial_K V > \ = \ < \nabla_K V > \ = 0.$ Here ∇ denotes the covariant derivate on the Kähler manifold defined by $\nabla_K V_M = \partial_K V_M - \Gamma_{KM}^L V_L$. Using the definition of the potential, eq.(1) and eq.(2), this implies[3]:

$$< e^G(G^M \ \nabla_K G_M + G_K) + \sum_A g_A^2 D_A(\partial_K D_A - \frac{1}{2}G_K D_A) >= 0 \tag{3}$$

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[1] Most of the expressions are presented in Planck units, namely, we set $M_P = 1$. However, at many instances, we keep M_P explicitly to make the discussions clearer.

[2] Note that the D-terms can be explicitly given in terms of the fields, derivatives of the Kähler potential and a FI term. We will make use of this form in a later subsection. For the present, we just note that we consider FI terms to be moduli dependent.

[3] Strictly speaking, there is a contribution proportional to the derivative of the gauge kinetic function in the minimisation condition, eq.(3). As we are concerned with the general expressions for the matter field soft terms, these contributions will be proportional to matter field vevs which are much smaller than the moduli *vevs* and therefore we will neglect them here.

We will use the eqs.(2,3) while deriving general expressions for the soft terms. In the present subsection we will not distinguish between the matter and hidden/moduli fields, but present generic expressions for the various scalar couplings in the theory. To start with, we will consider the case of the scalar mass squared matrix, which is defined as

$$M_0^2 = \begin{vmatrix} m_{I\bar{J}}^2 & m_{IJ}^2 \\ m_{\bar{I}\bar{J}}^2 & m_{\bar{I}J}^2 \end{vmatrix}, \tag{4}$$

where the various entries are defined by:

$$m_{I\bar{J}}^2 = \; <\partial_I \partial_{\bar{J}} V> \; = \; <\nabla_I \nabla_{\bar{J}} V> \tag{5}$$

$$m_{IJ}^2 = \; <\partial_I \partial_J V> \; = \; <\nabla_I \nabla_J V> . \tag{6}$$

Using the definition of the potential in eq.(1) and the conditions, eqs.(2, 3), we find the most general expressions for the bilinear couplings to be of the form :

$$m_{I\bar{J}}^2 = e^G(G_{I\bar{J}} + \nabla_I G^{\bar{K}} \nabla_{\bar{J}} G_{\bar{K}} - R_{I\bar{J}K\bar{L}} G^K G^{\bar{L}}) + \frac{1}{2}\sum_A g_A^2 D_A^2 (G_{\bar{J}} G_I - G_{I\bar{J}})$$
$$- \sum_A g_A^2 D_A (G_{\bar{J}} \partial_I D_A + G_I \partial_{\bar{J}} D_A - \partial_I \partial_{\bar{J}} D_A) + \sum_A g_A^2 \partial_I D_A \partial_{\bar{J}} D_A , \tag{7}$$

$$m_{IJ}^2 = e^G(2\nabla_J G_I + G^K \nabla_I \nabla_J G_K) - \sum_A g_A^2 D_A (G_J \partial_I D_A + G_I \partial_J D_A - \nabla_I \nabla_J D_A)$$
$$- \frac{1}{2}\sum_A g_A^2 D_A^2 (G_I G_J + \nabla_I G_J + \frac{1}{2} g_A^2 \partial_I \partial_J f_A) + \sum_A g_A^2 \partial_I D_A \partial_J D_A , \tag{8}$$

where we have neglected the vacuum brackets for simplicity[4]. Here, f_A represents the gauge kinetic function. The term containing the second derivative $\partial_I \partial_J f_A$ gives contributions to the B_μ term from operators of the form $\int d^2\theta \, W^\alpha W_\alpha H_1 H_2$ in superfields. A similar term of the form, $\partial_I f_A \partial_{\bar{J}} \bar{f}_A$ could contribute to $m_{I\bar{J}}^2$. However since this contribution is proportional to the vevs of the matter fields, as we will discuss in the next section, we neglect it here[5]. The function, $R_{I\bar{J}K\bar{L}}$ represents the Riemann (curvature) tensor of the Kähler manifold whose definition can be found in any of the standard texts [13]. The next step would be to derive the expression for the trilinear couplings, which we define as[6]:

$$A_{IJK} = <\nabla_I \nabla_J \nabla_K V> \tag{9}$$

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[4] From now on we will neglect vacuum brackets in the rest of the paper, unless and otherwise specified.

[5] For the same reason, we do not write down the contributions from gauge kinetic function in the A_{ijk} term discussed below.

[6] For MSSM fields this definition is equivalent to the naive one of using ordinary derivatives giving the A-term.

which takes the form :

$$
\begin{aligned}
A_{IJK} \;=\; & e^G \Big(G_K(2\nabla_J G_I + G^M \nabla_I \nabla_J G_M) + G_J(2\nabla_K G_I + G^M \nabla_I \nabla_K G_M) \\
& + \; G_I(2\nabla_K G_J + G^M \nabla_J \nabla_K G_M) + 2\nabla_I \nabla_K G_J + \nabla_J \nabla_K G_I + G^M \nabla_I \nabla_J \nabla_K G_M \Big) \\
& - \; g_A^2 D_A (\nabla_I D_A - \tfrac{1}{2} G_I D_A)(G_J G_K + \nabla_J G_K) - g_A^2 D_A (\nabla_J D_A - \tfrac{1}{2} G_J D_A)(G_I G_K + \nabla_I G_K) \\
& - \; g_A^2 D_A (\nabla_K D_A - \tfrac{1}{2} G_K D_A)(G_I G_J + \nabla_I G_J) - \tfrac{1}{2} g_A^2 D_A^2 \, (G_I G_J G_K + G_I \nabla_J G_K \\
& + \; G_K \nabla_I G_J + G_J \nabla_I G_K + \nabla_I \nabla_J G_K) + g_A^2 \Big(\nabla_I \nabla_J D_A \nabla_K D_A + \nabla_J D^A \nabla_I \nabla_K D_A \\
& + \; \nabla_I D_A \nabla_J \nabla_K D_A + D_A \nabla_I \nabla_J \nabla_K D_A \Big) \,.
\end{aligned}
\tag{10}
$$

Note that as for the B term, there can be contributions to the A-term also from gauge kinetic function, which can be represented by operators of type $\int d^2\theta W^\alpha W_\alpha h_{ijk} Q_i Q_j Q_k$ with Q_i representing the matter fields. These are typically of the order $m_{3/2}^2/M_{Pl}$ and thus they give negligibly small contributions unless $m_{3/2}$ has intermediate scale values. In the next subsection, we will use these expressions to get the expressions of soft masses for the matter fields.

A. General Expressions of soft terms for Matter Fields

We will define matter fields by setting their *vev*s to zero. This would mean that both the F and D contributions proportional matter field vevs to be zero at the leading order. Thus, we have:

$$
\langle \Phi^i \rangle = 0 \quad , \quad \langle G^i \rangle = 0 \quad , \quad \langle \partial_i D_A \rangle = 0 \quad ,
$$

with Φ representing the scalar part of a matter field. From now on, to distinguish matter and hidden/moduli fields, we denote matter (moduli/hidden sector) fields by using latin(greek) indices. To derive the soft terms for matter fields from the general scalar couplings presented in the previous section, along with using the definitions above, we have to remove the supersymmetric contributions from them. Further, we identify the gravitino mass to be $m_{3/2} = < e^{G/2} >$. Taking all these modifications in consideration, the final set of equations are of the form:

$$
m_{i\bar{j}}^2 \;=\; m_{3/2}^2 \, (G_{i\bar{j}} - R_{i\bar{j}\alpha\bar{\beta}} G^\alpha G^{\bar{\beta}}) - \frac{1}{2} \sum_A g_A^2 D_A^2 G_{i\bar{j}} + \sum_A g_A^2 D_A \partial_i \partial_{\bar{j}} D_A \;,
\tag{11}
$$

$$
m_{ij}^2 \;=\; m_{3/2}^2 \, (2\nabla_i G_j + G^\alpha \nabla_i \nabla_j G_\alpha) - \frac{1}{2} \sum_A g_A^2 D_A^2 (\nabla_i G_j + \frac{g_A^2}{2} \partial_i \partial_j f_A) + \sum_A g_A^2 D_A \nabla_i \nabla_j D_A
\tag{12}
$$

$$
A_{ijk} \;=\; m_{3/2}^2 \, (3\nabla_i \nabla_j G_k + G^\alpha \nabla_i \nabla_j \nabla_k G_\alpha) - \frac{1}{2} \sum_A g_A^2 D_A^2 \nabla_i \nabla_j G_k + \sum_A g_A^2 D_A \nabla_i \nabla_j \nabla_k D_A \;,
\tag{13}
$$

$$
\mu_{ij} \;=\; m_{3/2} \, \nabla_i G_j \quad , \quad M_{1/2}^A = \frac{1}{2}(Re f_A)^{-1} m_{3/2} f_{A\alpha} G^\alpha \;,
\tag{14}
$$

where we have now also supplemented the scalar equations with those for the μ and the gaugino masses. In the above $f_{A\alpha} = \partial f_A / \partial z^\alpha$. Note that these expression reduces to the standard form[5, 13, 14] in the limit where D_A goes to zero.

Note that the above soft terms are not in a canonically normalised basis for the kinetic terms. This can be seen from their action which has the form $g_{i\bar{j}} \partial z^i \partial z^{\bar{j}} - m^2_{i\bar{j}} z^i z^{\bar{j}}$, where z^i represents a matter scalar field. To go to the normalised basis, one can define vielbeins such as : $z^i = e^i_a z^a$, $z^{\bar{j}} = e^{\bar{j}}_{\bar{b}} z^{\bar{b}}$ such that $e^i_a e^{\bar{j}}_{\bar{b}} = g^{i\bar{j}}$. Using this transformations, we have the soft mass in the normalised basis to be given by

$$\bar{m}^2_{a\bar{b}} = e^i_a m^2_{i\bar{j}} e^{\bar{j}}_{\bar{b}},$$

where $\bar{m}^2_{a\bar{b}}$ represents the normalised masses. Similar analysis can be extended for other soft terms. In order to keep a compact notation, we however do not present the general expressions in the normalised form.

While these expressions are given for the tree level potential, higher order corrections can play a significant role, depending on the specifics of the model of supersymmetry breaking. In models with small tree-level contributions, the dominant set of corrections are of anomaly mediated type[15] which are proportional to the gravitino mass $m_{3/2}$. These contributions are typically not modified in the presence of D-terms and have to be anyway included. For the gauginos, the most general form of these expressions have been presented in [16] and are given by

$$M'^A_{1/2} = -\frac{g^2_A}{16\pi^2} \left(3T^A_G - T^A_R - (T^A_G - T^A_R) K_\alpha G^\alpha - \frac{2T^A_R}{d^A_R} (\log \det K|^"_R)_{,\alpha} G^\alpha \right) m_{3/2} . \qquad (15)$$

Here, T_G is the Dynkin index of the adjoint representation, normalised to N for $SU(N)$, T_R is the Dynkin index associated with the representation R of dimension d_R, normalised to $1/2$ for the fundamental of $SU(N)$ and $K|^"_R$ is the Kähler metric restricted to the representation R. This expression reduces to the following when all the *vev*s are much less than M_P:

$$M'^A_{1/2} = -\frac{g^2_A b^A_0}{16\pi^2} m_{3/2} , \qquad (16)$$

where the beta function b^A_0 was given as $3T^A_G - T^A_R$ in the previous expression. In addition to the gauginos, the scalar mass terms as well as the B-term and the A-terms receive corrections. In the case of gauginos, as long as the tree-level F-term contributions are present, the anomaly mediated contributions remain sub-dominant, whereas in the case of scalar soft terms, both the D-term as well as the F-term contributions have to be suppressed for the anomaly mediated contributions to dominate.

B. Implications of large D-terms on the soft parameters

Eqs.(11-14) give the modified expressions for the soft terms after including non-zero D-type SUSY breaking contributions in supergravity. Whereas the scalar couplings receive corrections from the D-type terms, the gaugino masses are unaffected by D-mediated effects. The μ term, could be visualised as a soft mass in supergravity by using the Giudice-Masiero mechanism[17]. The expression presented in the previous sub-section takes care of this situation and it is seen that D terms do not effect the μ term either. However, the exact implications on the soft terms by the inclusion of the D-terms depend on (a) the structure of the D-terms and (b) the magnitude of them. We will address these two issues below.

In the presence of anomalous non-linearly realised abelian gauge symmetries

$$\delta V_A = \Lambda_A + \bar{\Lambda}_A \ , \ \delta z^i = \Lambda_A X_A^i z^i \equiv \mathcal{V}_A^i \Lambda_A \ ,$$
$$\delta T^\alpha = \eta_A^\alpha \Lambda_A \equiv \mathcal{V}_A^\alpha \Lambda_A \ , \tag{17}$$

where $\mathcal{V}_A^I$ are the Killing potentials, the auxiliary D-terms, defined by

$$\partial_{\bar{J}} D_A = \mathcal{V}_A^I K_{\bar{J},I} = \mathcal{V}_A^i K_{\bar{J},i} + \mathcal{V}_A^\alpha K_{\bar{J},\alpha} \tag{18}$$

are explicitly given by

$$D_A = z^I X_I^A \frac{\partial K}{\partial z^I} + \xi_A = \bar{z}^{\bar{I}} X_I^A \frac{\partial K}{\partial \bar{z}^{\bar{I}}} + \xi_A \ , \ \xi_A \equiv \eta_A^\alpha \ \partial_\alpha K \ , \tag{19}$$

where X_I^A represents the $U(1)_A$ charges of the fields z^I and ξ_A denotes the Fayet-Iliopoulos term for the $U(1)_A$ factors. Note that the equality between the two last terms is a straightforward consequence of the gauge invariance of the Kähler potential. We consider the Fayet-Iliopoulos terms to be moduli dependent and we will not explicitly discuss here the various possible mechanisms of moduli stabilisation[7]. We have in the vacuum, after setting the matter fields *vevs* to zero

$$\langle \partial_j D_A \rangle = \langle \bar{v}_{\bar{\beta}} X_{\bar{\beta}}^A K_{\bar{j}\beta} + \eta_A^{\bar{\alpha}} K_{j\bar{\alpha}} \rangle = 0 \quad , \quad \langle \nabla_i \nabla_j D_A \rangle = 0 \ ,$$
$$\langle \partial_i \partial_{\bar{j}} D_A \rangle = K_{i\bar{j}} X_i^A + (\bar{v}^{\bar{l}} X_A^{\bar{l}} \partial_{\bar{l}} + \eta_A^{\bar{\alpha}} \partial_{\bar{\alpha}}) \ K_{i\bar{j}} \ , \ \langle \nabla_i \nabla_j \nabla_l D_A \rangle = 0 \tag{20}$$

By using (20), the soft terms for the matter fields reduce to

$$m_{i\bar{j}}^2 = m_{3/2}^2 \left(G_{i\bar{j}} - R_{i\bar{j}\alpha\bar{\beta}} G^\alpha G^{\bar{\beta}} \right) + \sum_A g_A^2 D_A \left(X_i^A + \bar{v}_{\bar{l}} X_{\bar{l}}^A \partial_{\bar{l}} + \eta_A^{\bar{\alpha}} \partial_{\bar{\alpha}} - \frac{1}{2} D_A \right) G_{i\bar{j}} \ , \tag{21}$$

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[7] After moduli stabilisation, the anomalous $U(1)$'s become gauged R-symmetries [18].

$$m_{ij}^2 = m_{3/2}^2 \ (2\nabla_i G_j + \ G^\alpha \nabla_i \nabla_j G_\alpha) - \frac{1}{2} \sum_A g_A^2 D_A^2 (\nabla_i G_j + \frac{g_A^2}{2} \partial_i \partial_j f_A) \ , \tag{22}$$

$$A_{ijk} = m_{3/2}^2 \ (3\nabla_i \nabla_j G_k + G^\alpha \nabla_i \nabla_j \nabla_k G_\alpha) - \frac{1}{2} \nabla_i \nabla_j G_k \sum_A g_A^2 D_A^2 \ . \tag{23}$$

Let us now try to quantify how large the D-terms can be. To do this, let us consider the following generic forms for the Kähler and the superpotential :

$$K = \tilde{K}(T_\alpha, T_{\bar\beta}) + H_{i\bar j}(T_\alpha, T_\beta) Q_i Q_{\bar j}^\dagger + (Z_{ij}(T_\alpha, T_{\bar\beta}) Q_i Q_j + h.c) + \dots \tag{24}$$

$$W = Y_{ijk}(T_\alpha) Q_i Q_j Q_k + \tilde{W}(T_\alpha) + \dots, \tag{25}$$

where T_α represent moduli/hidden sector fields and Q_i represent the matter fields. Using these equations let us now revisit the condition (2)

$$m_{3/2}^2 \left(K^{\alpha\bar\beta}(K_\alpha K_{\bar\beta} + \frac{M_P^2}{W}(K_\alpha W_{\bar\beta} + W_\alpha K_{\bar\beta}) + \frac{M_P^4}{|W|^2} W_\alpha W_{\bar\beta}) - 3M_P^2 \right) + \frac{1}{2} \ g_A^2 D_A^2 = 0 \ . \tag{26}$$

From the above we see that, as long as the D-terms are in the limit, $D \sim \mathcal{O}(m_{3/2}^2)$, they would not contribute significantly to the vacuum energy. However, when they lie within the limit

$$m_{3/2}^2 \lesssim D_A \lesssim m_{3/2} M_P, \tag{27}$$

they could be contributing significantly. The upper limit is obtained when one assumes D-term contributions to dominate over the F-term contributions or are of the same order as them. This particular limit is what we are interested in the present work as this has not been exploited in a general manner as presented here. From the generic set of soft parameters presented above, it is obvious that splitting between fermionic and scalar superpartners can be 'naturally' achieved once the D-terms lie within the above range. Quantitatively, if in a given model the gravitino mass is of $\mathcal{O}(1 \text{ TeV})$, the upper limit on the D-term would be of the order of intermediate scale $\sim (10^{10})$ GeV. It is obvious that as one increases the gravitino mass closer to the intermediate scale $\sim (10^9 - 10^{12})$ GeV, the upper bound on the D-terms become close to the GUT scale. These upper bounds are essentially the magnitude required to cancel the cosmological constant in the limit where the F-terms tend to zero.

Given this limit, let us now try to understand in more detail how large D-terms would generate large splittings between superpartners. The equations for the gaugino and μ-term remain unchanged as we have mentioned. The following features of the spectra are easy to extract without actually being specific about the model:

- (i). *Scalar Mass Terms*: The most dominant contribution to the scalar masses from the D-terms are the ones which are linear in D which for $m_{3/2} \sim$ TeV push the scalar masses

to intermediate energy scale. Note that these terms depend on the charges of the fields under the additional $U(1)$ gauge group, thus putting a constraint that these charges to be of definite sign. If all the three generations of the sfermions have the same charges under the $U(1)$ groups, this term would also be universal. Otherwise, there are off-diagonal entries which are generated in the mass matrices, which could of suppressed by some powers in the expansion parameter $\epsilon_\beta = v_\beta/M_P$, with v_β representing the *vev* of some flavon field.

- (ii). *Higgs mass terms and the B_μ*: The Higgs masses follow almost the same requirements as the soft masses. Usually, their charges are linked with the Giudice-Masiero mechanism[17]. The B_μ term is however special. Unlike the Higgs mass terms, it does not receive large contributions from D-terms, whose contributions can be utmost of $\mathcal{O}(m_{3/2}^2)$. If the splitting between the Higgs masses and the B_μ is too large, it could lead to unphysical regions in $\tan\beta$. This could be easily seen by noting that

$$\sin 2\beta = \frac{2B_\mu}{m_{H_1}^2 + m_{H_2}^2 + 2\mu^2}. \tag{28}$$

In the limit of large Higgs mass parameters $m_{H_1}^2$, $m_{H_2}^2$, one has to think of ways to enhance the B_μ term. We will present one such example in the next section.

- (iii). *A-terms*: Even if the D-terms are large, the A-terms are typically proportional to $\mathcal{O}(m_{3/2})$. No large enhancement is present. This is expected as A-terms break R-symmetries. They get related to the D-terms due to the constraints of cosmological constant cancellation, but as the scale of R-symmetry breaking is set by the gravitino mass, this naturally sets the A-terms to be of same order.

- (iv). *Gaugino Masses*: All along we have been commenting that the presence of SUSY breaking D-terms would not change the results for the gaugino masses presented there. This is only true as long as there are no additional fermions in the model. In the presence of additional fermions and non-zero D-terms, gauginos can get Dirac masses through operators of the form [23]

$$h_a \int d^2\theta \, \frac{\chi^a W_\alpha^a W_X}{M_P l} \;\; = \;\; h_a \frac{\langle D_X \rangle}{M_P} \psi^a \lambda^a \;\; + \cdots \;\; = \;\; m_D^a \psi^a \lambda^a + \cdots \,, \tag{29}$$

where χ^a represent here fields in the adjoint representation of the Standard Model gauge group with (mirror fermions) which mix with the gauginos and m_D^a represent the Dirac mass for the gauginos. These mixing terms could lead to the Majorana masses for the gauginos by

a seesaw mechanism $\sim (m_D^a)^2/M_a$ if the mirror fermions obtain large R-symmetry breaking Majorana masses, M_a. In the present work, we do not concentrate on building models of this type.

III.   PARAMETRIZATION OF SOFT TERMS IN TYPE I/II STRING MODELS WITH LARGE D-TERMS

The soft terms in effective string supergravities from Type-I/II string theories have been parameterized in [6] where pure F-type breaking has been assumed. In the present section we will extend this analysis by considering D-type SUSY breaking terms too. In each of this case, we present parameterizations of the soft terms which could be readily be useful for phenomenological studies.

A.   D-dominated supersymmetry breaking

The first case we consider is that of a scenario where F-terms are absent or negligible. We assume that supersymmetry breaking is achieved by pure D-terms. However, we will still require that the gravitino get a mass. This would enable us to cancel the cosmological constant even in the pure D-breaking limit[8]. The scale of the gravitino mass is assumed to be not very far from the weak scale. With these conditions, the potential, eq.(1) takes the form:

$$V = \frac{1}{2}\sum_A g_A^2 D_A^2 - 3m_{3/2}^2 M_P^2 \ . \tag{30}$$

It is obvious from the above equation that requiring that the potential should vanish at the minimum (for the cosmological constant cancellation), implies that the D-terms should be

$$<D> = \frac{\sqrt{6}}{g}\ m_{3/2} M_P \ . \tag{31}$$

A more subtler constraint comes from the existence of a minimum, eq.(3). In this limit, it takes the form $g_A^2 D_A(\partial_\beta D_A) = 0$. It is clear that for a single $U(1)$ gauge group, this would mean at the minimum either the *vev* to vanish or the D-term to vanish. Both these conditions are not acceptable to us. The situation would not change even if one adds more flavon fields. Thus we rule out the case of single $U(1)$ with pure D-breaking. The minimum case we can think of is that case with two $U(1)$ gauge groups with two charged fields.

---

[8] An earlier proposal in this direction has been presented in [25].

We parameterise the SUSY breaking D-terms, consistently with the vanishing of the cosmological constant, as

$$< D_A > = \frac{\sqrt{6}}{g_A} \, \theta_A \, m_{3/2} M_P \; , \tag{32}$$

where θ_A are defined such that $\sum_A \theta_A^2 = 1$. Then the soft terms reduce to the following form :

$$
\begin{aligned}
m_{i\bar{j}}^2 &= -2m_{3/2}^2 \, G_{i\bar{j}} + \sqrt{6} \, m_{3/2} M_P \sum_A g_A \theta_A (X_i^A + v_{\bar{l}} X_{\bar{l}}^A \partial_{\bar{l}} + \eta_A^{\bar{\alpha}} \partial_{\bar{\alpha}}) G_{i\bar{j}} \; , \\
m_{ij}^2 &= -m_{3/2}^2 \left(\nabla_i G_j + \frac{3}{2} \sum_A g_A^2 \theta_A^2 \partial_i \partial_j f_A \right) \; , \\
\mu_{ij} &= m_{3/2} \, \nabla_i G_j \; , \\
A_{ijk}' &= m_{3/2} \lambda_{ijk} (\gamma_i + \gamma_j + \gamma_k) \; , \\
M_{1/2}^{'\,A} &= -\frac{g_A^2 b_0^A}{16\pi^2} \, m_{3/2} \; .
\end{aligned}
\tag{33}
$$

The gaugino masses vanish at the tree level in this limit. They are generated by anomaly mediated contributions as listed above. Similar thing happens for the A-parameters, which are determined by their anomalous dimensions (γ_i) as given above. Note that the above mass formulae are given at the high scale. One has to evolve these masses at the weak scale to make contact with weak scale phenomenology. The present scenario describes a new situation where the non-holomorphic scalar soft masses are given by dominant D-type supersymmetry breaking terms, whereas the gauginos, described by the beta-functions, the supersymmetric fermion masses (in particular the μ term of MSSM) are proportional to the gravitino mass and have therefore much lower values. If all the $U(1)$ groups are in the visible sector with large D-terms and positive charges, such a situation is not phenomenologically viable, since there is no possibility of tuning one Higgs doublet to be very light. However, if some of the $U(1)$ lie in the hidden sector with some others in the visible sector and the angles θ_A in the visible sector are all small, then the scenario with pure D-breaking becomes viable. In this last case, all the soft terms can be at the TeV scale, thus making contact with a low energy physics of the MSSM type. It would be interesting to see how this new structure of soft terms would feature with respect to low-energy constraints like electroweak symmetry breaking, dark matter, LEP Higgs bounds and other constraints. Note that a situation like split SUSY could be difficult to incorporate here.

B.   D-breaking with dilaton and moduli supersymmetry breaking

The above discussion presents an extreme situation *i.e.* completely absent F-type breaking. However such an extreme limit is not required to realise split supersymmetry breaking. The

general analysis presented in the previous section shows that it is enough to have $g_A^2 D_A >> m_{3/2}^2$. We present here soft terms for a case where, for simplicity, there is only one $U(1)$ large D-term and we assume that the auxiliary field of the dilaton or the overall modulus superfields also contribute to supersymmetry breaking.

Note that such a situation can arise naturally when one considers effective lagrangians of Type I string theory for an orientifold with only $D9$ branes. We provide expressions for the case of orbifold theories (Calabi-Yau spaces are also particular cases of the expressions below) in the large volume limit. In this limit, the gauge kinetic function and the Kähler potential K will have the general form [27]:

$$
\begin{aligned}
f_A^B &= S\,\delta_A^B + \dots\ , \\
K &= -\log(S + S^\dagger) - 3\log(T + T^\dagger - \delta_{GS}V) + \sum_i (T + T^\dagger - \delta_{GS}V)^{n_i}|\phi_i|^2 \\
&\quad + \sum_{ijk}\left(Z_{ijk}(T + T^\dagger - \delta_{GS}V)^{n_i}\bar{\phi}_i\phi_j\phi_k + \text{h.c.}\right) + \cdots\ ,
\end{aligned}
\tag{34}
$$

where we have used the by now standard notation with S representing the dilaton field, T representing the overall volume modulus, ϕ_i represent matter fields and n_i modular weights of matter fields. We are assuming from now on that the modulus T is the one mixing with the anomalous $U(1)$ gauge field, such that the gauge invariant combination $T + T^\dagger - \delta_{GS}V$ should consistently appear in the Kähler potential and in the couplings to the matter fields. This can be explicitly realized in intersecting brane models, as we will illustrate later on. The last term in the Kähler potential in (34) accommodate the possibility of μ terms and simultaneously, that of the B_μ term. We parameterize the supersymmetry breaking contributions from the two sets of auxiliary fields as :

$$
< G_S >= \sqrt{3}\,\left(\frac{M_P}{S + S^\dagger}\right)\,\cos\gamma_S\quad,\quad < D >= \frac{\sqrt{6}}{g}m_{3/2}M_P\sin\gamma_S\ .
\tag{35}
$$

We then obtain the soft terms

$$
\begin{aligned}
m_{i\bar{j}}^2 &= (1 - 3\sin^2\gamma_S)m_{3/2}^2 G_{i\bar{j}} + \sqrt{6}\,g\,m_{3/2}M_P\sin\gamma_S(X_i + \bar{v}_{\bar{l}}X_{\bar{l}}\partial_{\bar{l}} + \delta_{GS}\partial_{\bar{T}})G_{i\bar{j}} \\
m_{ij}^2 &= (2 - 3\sin^2\gamma_S)m_{3/2}^2\nabla_i G_j\ -\frac{3}{2}m_{3/2}^2 g^2\sin^2\gamma_S\ \partial_i\partial_j f\ , \\
A_{ijk} &= 3m_{3/2}^2\cos^2\gamma_S\nabla_i\nabla_j G_k\ , \\
M_{1/2}^A &= \frac{\sqrt{3}}{2}m_{3/2}\cos\gamma_S\ ,
\end{aligned}
\tag{36}
$$

whereas the μ term is unchanged (33). In the complementary case where the only F-type source of supersymmetry breaking comes from the T field, the appropriate parametrization is

$$< G_T >= -3 \left(\frac{M_P}{T + T^\dagger}\right) \cos \gamma_T \quad , \quad < D >= \frac{\sqrt{6}}{g} m_{3/2} M_P \sin \gamma_T . \tag{37}$$

The soft terms in this case are given by

$$
\begin{aligned}
m_{i\bar{j}}^2 &= (1 + n_i \cos^2 \gamma_T - 3 \sin^2 \gamma_T) m_{3/2}^2 G_{i\bar{j}} + \sqrt{6}\, g\, m_{3/2} M_P \sin \gamma_T (X_i + \bar{v}_{\bar{l}} X_{\bar{l}} \partial_{\bar{l}} + \delta_{GS} \partial_{\bar{T}}) G_{i\bar{j}}, \\
m_{ij}^2 &= [2 + (n_i + n_j) \cos \gamma_T - 3 \sin^2 \gamma_T] m_{3/2}^2 \nabla_i G_j - \frac{3}{2} m_{3/2}^2 g^2 \sin^2 \gamma_T\ \partial_i \partial_j f\ , \\
A_{ijk} &= m_{3/2}^2\ [3 \cos^2 \gamma_T + (n_i + n_j + n_k) \cos \gamma_T]\ \nabla_i \nabla_j G_k\ , \\
M'^A_{1/2} &= -\frac{g_A^2}{8\pi^2} \left(3 T_G^A \sin^2 \frac{\gamma_T}{2} - T_R^A (\sin^2 \frac{\gamma_T}{2} - (1 + n_i) \cos \gamma_T) \right)\ m_{3/2}\ .
\end{aligned}
\tag{38}
$$

Several simplifying assumptions were used in deriving (38). For reasons already explained, the analytic scalar masses come from a Giudice-Masiero term in the Kähler potential of the type $\phi^\dagger Q_i Q_j + $h.c., where ϕ is a flavon type field with a large vev. The Yukawa couplings were assumed, in the large volume limit, to become T-modulus independent, otherwise new contributions appear in the trilinear A-terms. The natural values of modular weights for charged D9 branes charged fields are $n_i = -1$. Finally for phenomenological studies, the angles $\gamma_{S,T}$ can be used as independent parameters to be constrained by low energy physics.

IV. SPLIT SUPERSYMMETRY

The requirement of split supersymmetry type soft spectra are as follows :

(i) Scalar soft terms : $m_{\tilde{f}}^2 \sim \mathcal{O}(10^6 - 10^{15})$ GeV, $(\tilde{f} = Q,\ u^c,\ d^c,\ L,\ e^c)$

(ii). Higgs mass parameters $m_{H_1}^2 \sim m_{H_2}^2 \sim B_\mu \sim \mathcal{O}(10^6 - 10^{15})$ GeV, with one of the Higgs mass eigenvalues fine tuned to be around the electroweak scale.

(iii). The gaugino masses and the μ term are around the weak scale.

As a starting point, let us consider for a moment that all D-term contributions are negligible or zero. In such a case, we see that most likely the mass squared terms are proportional to $m_{3/2}^2$ whereas the gaugino masses are proportional to $m_{3/2}$. Thus, it is difficult to expect a large splitting within the masses of the superpartners in supergravity theories with pure or dominant F-type SUSY breaking. In principle, such a splitting can be arranged by choosing suitable parameter space within the goldstino directions in certain classes of effective lagrangians coming from heterotic strings. However, it is not clear how much these parameter spaces would remain stable

under radiative corrections. Another approach for creating a split would be to assume some R-symmetries[9] protecting the fermion superpartners. In this case, the gravitino mass needs to be pushed to very high values, whereas the gauginos need another mechanism to achieve masses close to the weak scale[10, 22]. However in this case, one has to invent a mechanism to suppress the anomaly mediated contributions, which could involve for example no-scale type models.

In the presence of D-terms, it is generically difficult to realise split supersymmetry like models [10]. From the discussion in the previous section, it was obvious that it is just not sufficient to choose the $U(1)$ charges of the scalars to be positive to realise the split spectrum since B_μ term does not have large D-term contributions, we need to disentangle the μ and the B_μ term by introducing a new field X and allowing a term of the type XH_1H_2 in the superpotential. In a simple example, the field content is as follows. The model contains an additional $U(1)$ group, with two additional fields X and ϕ with charges $+2$ and -1. The ϕ field can act as a flavon field attaining a large vev close to the fundamental scale. The superpotential and the relevant term in the Kähler potential are specified as

$$W = W_{SSM} + \lambda_1 X H_1 H_2 + \lambda_2 X \phi^2 + \cdots ,$$

$$K \supset \sum_i |\phi_i|^2 + (\phi^\dagger)^2 H_1 H_2 + \cdots . \tag{39}$$

The scalar potential at the global SUSY level is given by

$$V = \lambda_2^2(|\phi|^4 + 4|X|^2|\phi|^2) + \frac{1}{2}g^2 \left(2|X|^2 - |\phi|^2 + \xi\right)^2 + \dots . \tag{40}$$

For $\xi > 0$, the stable extremum of the above potential and the auxiliary fields are given by:

$$\langle \phi \rangle = \frac{g^2}{2\lambda_2^2 + g^2}\xi , \ \langle X \rangle = 0,$$

$$\langle F_\phi \rangle = 0 , \ \langle F_X \rangle = \frac{\lambda_2 g^2}{2\lambda_2^2 + g^2}\xi , \ \langle D \rangle = \frac{2\lambda_2^2}{2\lambda_2^2 + g^2}\xi . \tag{41}$$

From the above it is clear that $F_X \sim g^2 D$ and moreover of the order of the FI term ξ. This is sufficient to enable the B term to receive large contributions through the term $G^X \nabla_{H_1} \nabla_{H_2} G_X$ in the eq.(22). As long as ξ is close to an intermediate scale value, this model seems to replicate the split spectrum, if one fixes the gravitino mass around 1 TeV. However, in typical string models, the FI term is of the $\mathcal{O}(M_{Pl}^2/16\pi^2)$ which would give a too large contribution to the vacuum energy. One way to get the correct order of magnitude is by incorporating the above model into a higher

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[9] Or even a charge symmetry accompanied by F-breaking of charged chiral superfield, [21].
[10] See also [26].

dimensional theory. For illustration lets us consider a 5D theory compactified over S^1/Z_2. The Standard Model and the X, ϕ fields live on a 3D brane, whereas the gauge fields of the $U(1)$ are allowed to propagate in the bulk. We will use Scherk-Schwarz mechanism to break supersymmetry. The R-symmetry is also broken by this mechanism giving rise to the gravitino mass.

The various scales in the problem are $R = tM_5^{-1}$, $RM_5^3 = M_P^2$, where $t \equiv Re\ T$, the modulus field. After canonically normalizing the various fields by $\hat{\phi}_i = \sqrt{t/3}\ \phi_i$ and at the global super-symmetry level, the potential retains the form (40) with $\xi \sim M_5^2 = M_P^2/t$. The four dimensional $U(1)$ gauge coupling is given by $g^2 = 1/t = 1/(RM_5)$, whereas the gravitino mass is given by $m_{3/2} = \omega/R$, where ω is a number of order one. The D-term contribution to the vacuum energy is then of the form

$$\langle V_D \rangle \sim g^2 M_5^4 \sim m_{3/2}^2 M_P^2 \ , \tag{42}$$

in the right order as required by the cancellation of the vacuum energy in supergravity and realisation of the split spectrum. If the no-scale structure is broken by the dynamics, the gauginos attain their masses through anomaly mediation and thus we choose the gravitino mass to be of the order of 100 TeV. The μ is generated by the Giudice-Masiero mechanism and is $\mu \sim (<\phi>/M_5)^2 m_{3/2}$. So, this model replicates the spectrum of the split supersymmetry at the weak scale using large D-terms of the intermediate scale and a 100 TeV massive gravitino.

In the light of above discussion, an important question is in which sense the light Higgs mass tuning is preferred over the tuning of another scalar mass. Tuning of squarks or slepton masses is best described in terms of alignment in the 3×3 flavor space. If sfermion mass matrices are very close to the diagonal, i.e. off-diagonal terms are very small compared to the diagonal ones, the tuning of a small mass eigenvalue is impossible, whereas the tuning becomes more and more likely for off-diagonal terms of the same order as the diagonal ones. In flavor models with a low energy supersymmetric spectrum, the alignment of the quark-squark and lepton-slepton mass matrices was necessary to avoid too large FCNC effects, but a serious tension between alignment and hierarchy of fermion masses was present, at least for models with only one $U(1)$ factor. It is ironical that, in the limit of evading FCNC effects by decoupling the undesirable scalar particles, the alignment has still to be invoked in order to minimize the likelihood of the fine-tuning of squark and slepton masses compared with the tuning of the light Higgs mass.

V. NONPERTURBATIVE MODULI STABILISATION AND LARGE D-TERMS

In string theory, the FI terms are field (moduli) dependent. If no additional dynamics is present, the moduli fields will always exhibit a runaway behaviour and the FI terms disappear. We revisit here the issue of moduli stabilisation with realisation of large D-term contributions in a context similar to, but having some new features compared to the one discussed some time ago in [8]. As will become transparent, our analysis is also relevant for the issue of the uplift of the energy density in the context of KKLT type moduli stabilisation [35, 36]. The gauge group consists of the Standard Model supplemented by a confining hidden sector group and an anomalous $U(1)_X$. We consider the case of a supersymmetric $SU(N_c)$ gauge group with N_f quark flavors Q_i^a and anti-quark $\tilde{Q}_i^a$ where $a = 1 \cdots N_c$ is an index in the fundamental representation of the $SU(N_c)$ gauge group and $i, \bar{i} = 1 \cdots N_f$ are flavor indices. In the intersecting string realisation, discussed in some detail in the next section, the hidden sector consists of a stack of N_c magnetised D9 branes in the type I string with kinetic function $f = S + kT$, where S is the dilaton (super)field, T a volume (Kähler) modulus and k is a positive or negative integer determined by the magnetic fluxes in two compact torii. The low energy dynamics is described by $M_{\bar{j}}^i = Q^{a,i} \tilde{Q}_{\bar{j}}^a$, the composite "mesons" fields. In the following we denote by q $(\bar{q})$ the $U(1)_X$ charges of the hidden sector quarks (antiquarks). Since the FI terms are T-modulus dependent, T will shift under gauge transformations

$$
\begin{aligned}
V_X &\rightarrow V_X + \Lambda_X + \bar{\Lambda}_X \ , \ M_{\bar{j}}^i \rightarrow e^{-2(q+\bar{q})\Lambda_X} \ M_{\bar{j}}^i \ , \\
T &\rightarrow T + \delta_{GS} \Lambda_X \ ,
\end{aligned}
\tag{43}
$$

where

$$
\delta_{GS} = \frac{C_{N_c}}{k} \ , \ C_{N_c} = \frac{1}{4\pi^2} N_f (q + \bar{q}) \ ,
\tag{44}
$$

is uniquely fixed by the requirement that the mixed $U(1)_X SU(N_c)^2$ anomaly, denoted C_{N_c} in (44), to be exactly canceled by the nonlinear transformation of ImT. Notice that the nonlinear transformation of T forces a chiral nature of the hidden sector with respect to the anomalous abelian gauge group, which in turn triggers supersymmetry breaking [8]. In order to be able to write gauge invariant mass terms for the mesons, a field with charge opposite in sign to the ones of the mesons has to be introduced, called ϕ in what follows, of charge -1 in our conventions. The dynamical scale of the hidden sector gauge group, the effective superpotential [30] and the Kähler potential are

$$
\Lambda = M_P \ e^{-8\pi^2 (S+kT)/(3N_c - N_f)} \ ,
$$

$$W = W_0(S) + (N_c - N_f) \left(\frac{\Lambda^{3N_c - N_f}}{det M} \right)^{\frac{1}{N_c - N_f}} + m_i^{\bar{j}} (\frac{\phi}{M_P})^{(q+\bar{q})} M_{\bar{j}}^i ,$$

$$K = -\ln (S + \bar{S}) - 3 \ln [T + \bar{T} - 2 Tr(M^\dagger M)^{1/2} - |\phi|^2 - \delta_{GS} V] , \qquad (45)$$

where $m_i^{\bar{j}}$ are mass parameters. Notice first of all that the dynamical superpotential

$$W_{np} = (N_c - N_f) \left(\frac{e^{-8\pi^2(S+kT)}}{det M} \right)^{\frac{1}{N_c - N_f}} , \qquad (46)$$

is precisely gauge invariant when the anomaly cancellation conditions (43)-(44) are satisfied. In order to stabilise the modulus S we invoke the three-form NS-NS and RR fluxes. W_0 depends on the modulus S, $S = S_0$ and eventually other (complex structure) moduli of the theory and stabilises them by giving them a very large mass. If the other relevant mass scales, the FI term and the dynamical scale Λ have much lower values, we can safely integrate out these fields, by keeping the T modulus in the low energy dynamics. The resulting lagrangian is similar to the one invoked in the KKLT moduli stabilisation [35] with a D-term uplifting of the vacuum energy [36]. Notice however that the simple nonperturbative superpotential e^{-aT} considered in [36] cannot be gauge invariant due to the gauge transformation of T and therefore, precisely as in the heterotic case discussed in [8], charged hidden sector matter with appropriate charges is crucial to define a consistent gauge invariant model.

Minimisation with respect to T in (45) stabilises also the Kähler modulus. For notational simplicity we discuss in some detail the case of an supersymmetric hidden sector $SU(2)$ gauge group with one quark flavor Q^a and anti-quark $\tilde{Q}^a$ where $a = 1, 2$ is an index in the fundamental representation of the gauge group. Due to the anomalous nature of the $U(1)_X$, the sum of the quark and antiquark charges, equal to the M meson charge, is different from zero and, in our example, equal to $+1$. ϕ is a field of charge -1 which participate in the Yukawa coupling $\lambda \phi Q^a \tilde{Q}^a$, which plays the role of meson mass after the spontaneous symmetry breaking of the $U(1)_X$. The fact that the meson masses come from a perturbative trilinear Yukawa coupling in this case is instrumental in producing a large D-term contribution to supersymmetry breaking. In order to provide explicitly the scalar potential, we define the canonical field $M \equiv \chi^2/2$. Then the supergravity scalar potential can be found to be

$$V_F = \frac{1}{r^3} \left\{ \frac{r^2}{3} |\partial_T W - \frac{3}{r} W|^2 + \frac{r}{3} \sum_{i=1}^2 |\partial_i W + \bar{\phi}_i \partial_T W|^2 - 3|W|^2 \right\} ,$$

$$V_D = \frac{1}{S + \bar{S} + k'(T + \bar{T})} \left(\frac{3}{r} X_i |\phi_i|^2 + 3 \frac{\delta_{GS} M_P^2}{T + \bar{T}} \right)^2 , \qquad (47)$$

where $\phi_i = \chi, \phi$, we have introduced $r \equiv (T + \bar{T} - \sum_i |\phi_i|^2)$, δ_{GS} represent the Green-Schwarz coefficient of the $U(1)_X$, and k' is the magnetic flux on the brane providing the anomalous $U(1)_X$.

By inserting (45) into (47) we find a model with all moduli stabilised. If we would ignore the $U(1)_X$ dynamics, for example, T would be stabilised as in [35] by solving $D_T W = 0$. In our case, the minimum $T_0 = \langle T \rangle$ will be shifted due to the D and new F contributions. A full supergravity analysis of the vacuum of (47) is possible but cumbersome. Due to this after stabilizing $S = S_0$ and $T = T_0$ by solving their equations of motion, we analyse the stabilisation of the other fields, for simplicity at the global supersymmetry level, as in [8], by a suitable rescaling of the fields and Yukawa coupling λ. For general N_c , N_f at the global level, the auxiliary fields and the scalar potential are

$$
(F^{\bar{M}})^{\bar{i}}_i = 2[(M^\dagger M)^{1/2}]^{\bar{i}}_{\bar{j}} \left[-(M^{-1})^{\bar{j}}_i \left(\frac{\Lambda^{3N_c - N_f}}{det M} \right)^{\frac{1}{N_c - N_f}} + m^{\bar{j}}_i (\frac{\phi}{M_P})^{(q+\bar{q})} \right] ,
$$

$$
\bar{F}_{\bar{\phi}} = \frac{q + \bar{q}}{M_P} (\frac{\phi}{M_P})^{q+\bar{q}-1} Tr(mM) ,
$$

$$
D_X = (q + \bar{q}) Tr(M^\dagger M)^{1/2} - |\phi|^2 + k\mu^2 ,
$$

$$
V = |F_\phi|^2 + \frac{1}{2}[(M^\dagger M)^{-1/2}]^{\bar{j}}_{\bar{i}} (F^{\bar{M}})^{\bar{i}}_i (F^M)^i_{\bar{j}} + \frac{g_X^2}{2} D_X^2 , \tag{48}
$$

where $\mu^2 = 3C_{N_c}/k^2(T + \bar{T})$ is a mass scale determined by the T-modulus vev. The new feature of (48) is that k and consequently the FI term can have both signs, whereas in the effective heterotic string framework worked out in [8], the FI term had only one possible sign.

In the limit $\Lambda << \mu$, the vacuum structure and the pattern of supersymmetry breaking in the two cases of k positive and negative are vastly different.

i) $k > 0$. In this case the vacuum can be determined as in [8], where it was analysed for arbitrary $N_f < N_c$ and arbitrary $q + \bar{q} > 0$ charges. Keeping one mass parameter $m^{\bar{j}}_i = m\delta^{\bar{j}}_i$, we find, to the lowest orders in the parameter ϵ defined by

$$
\epsilon \equiv \frac{M_0}{k\mu^2} = (\frac{\Lambda}{\sqrt{k}\mu})^{\frac{3N_c - N_f}{N_c}} \left[\frac{m}{M_P}(\frac{\sqrt{k}\mu}{M_P})^{q+\bar{q}-1} \right]^{\frac{N_f - N_c}{N_c}} , \tag{49}
$$

a hierarchically small scale of supersymmetry breaking

$$
\langle |\phi|^2 \rangle = k\mu^2 \left[1 + \epsilon N_f(q + \bar{q}) \right] , \ \langle M \rangle = M_0 \left[1 - \epsilon(q + \bar{q})^2 \frac{N_f(N_c - N_f)(2N_c - N_f)}{2N_c^2} \right] ,
$$

$$
g_X^2 \langle D_X \rangle = -\epsilon^2 \hat{m}^2 N_f^2 (q + \bar{q})^2 \left[1 - \frac{N_f}{N_c}(q + \bar{q}) \right] ,
$$

$$
\langle F_\phi \rangle = \epsilon \hat{m} \sqrt{k}\mu N_f(q + \bar{q}) , \ \langle F^{\bar{M}} \rangle = K^{M\bar{M}} \partial_M W = -\epsilon^2 \hat{m} k\mu^2 \frac{N_f(N_c - N_f)}{N_c}(q + \bar{q})^2 , \tag{50}
$$

where $\hat{m} \equiv m(\sqrt{k}\mu/M_P)^{q+\bar{q}}$.

ii) $k < 0$. Here we specifically consider the case $N_c = 2$, $N_f = 1$ and $q + \bar{q} = 1$. In this case we find, to the lowest order in the parameter $\epsilon' = [(g^2 + 2\lambda^2)^4/8\lambda^2 g^{10}](\Lambda^2/|k|\mu^2)^5$, a large scale of supersymmetry breaking

$$\langle \phi \rangle = \frac{(g^2 + 2\lambda^2)^2}{2\lambda g^4} \frac{\Lambda^5}{k^2 \mu^4} \left[1 + 3(-g^2 + 14\lambda^2)\epsilon' \right] \ , \ \langle M \rangle = -\frac{g^2}{g^2 + 2\lambda^2} k\mu^2 \left[1 - 2(g^2 + 14\lambda^2)\epsilon' \right] \ ,$$

$$\langle D_X \rangle \simeq \frac{2\lambda^2}{g^2 + 2\lambda^2} k\mu^2 \ , \ \langle F_\phi \rangle \simeq -\frac{\lambda g^2}{g^2 + 2\lambda^2} k\mu^2 \ , \ \langle F^{\bar{M}} \rangle \simeq \frac{g^2 + 2\lambda^2}{g^2} \frac{\Lambda^5}{k\mu^2 M_P^2} \ . \tag{51}$$

Interestingly enough, this second case generate a large scale for supersymmetry breaking with large F_ϕ and D-term contributions. At first sight, a breaking of supersymmetry at a scale larger than the dynamical scale Λ destroys the supersymmetric confining dynamics underlying the nonperturbative superpotential in (46). However, the breaking of supersymmetry in the hidden sector is described by the mass splitting in the "mesonic" sector, measured by the auxiliary field F_M. Its value in case ii) is very small and actually the same as in case i), suggesting that the confining dynamics is still essentially supersymmetric. In the case $q + \bar{q} > 1$ we expect the D-term contribution to have further suppressions since the mesons masses come now from a higher dimensional operator. Within this context, we expect our general analysis of D-term contributions to supersymmetry breaking to be of relevance for further studies of phenomenological models incorporating moduli stabilisation [37]. In the following section we describe string theory realisations based on intersecting brane models leading precisely to the case $q + \bar{q} = 1$.

VI.   INTERSECTING BRANE STRING REALISATION OF LARGE D-TERM SUPERSYMMETRY BREAKING

Even if reasonable from a supergravity point of view, it is not obvious that a large D-term supersymmetry breaking in string theory is possible. Indeed, it is well known from the heterotic string constructions that the presence of Fayet-Iliopoulos terms triggers vev's for charged fields which break the gauge symmetry rather than supersymmetry [28]. This can presumably be understood by noticing that the FI terms in the heterotic string arise at one-loop and therefore, if they would break supersymmetry, they would be a radiative breaking of supersymmetry which is known to be very hard to obtain [29]. It was suggested in [8] that at the nonperturbative level, gaugino condensation in the hidden sector in the presence of an anomalous $U(1)$ symmetry can break supersymmetry. However as in the previous section for case i) introduced there, the induced D-terms are of the order (or slightly larger) than the F^2/M_P type terms and cannot provide the large contributions we are advocating in this paper.

In the Type I or Type II strings, on the other hand, the FI terms appear generically at tree-level and we can expect the tree-level supersymmetry breaking to be possible with large D-terms. As we will see, this will realise case ii) discussed in the previous section. We present here an explicit example suggesting this is indeed possible, in the context of intersecting branes Type I orientifold models or, T-dual equivalently, with internal magnetic fields [31, 32]. We discuss also various ingredients such that supersymmetry breaking to be really possible. We consider an explicit example, even if it is clear that a large class of similar models can be constructed. The model is based on the $Z_2 \times Z_2$ Type I orbifold without discrete torsion with internal magnetic fields $H_i^{(a)} = (m_i^{(a)}/v_i n_i^{(a)})$ in the torus T^i, where v_i are the volumes of the three torii. The model contains four stacks of D9 branes, each stack containing four coincident branes. Three of the stacks are magnetised and the fourth one is non-magnetised, with wrapping numbers $(m_i^{(a)}, n_i^{(a)})$ equal to

$$
\begin{aligned}
M_1 &: (m_i^{(3)}, n_i^{(3)}) &=& (0,1) &,& (2,1) &,& (2,1) \,, \\
M_2 &: (m_i^{(1)}, n_i^{(1)}) &=& (2,1) &,& (0,1) &,& (2,1) \,, \\
M_3 &: (m_i^{(2)}, n_i^{(2)}) &=& (-2,1) &,& (2,1) &,& (0,1) \,, \\
M_4 &: (m_i^{(4)}, n_i^{(4)}) &=& (0,1) &,& (0,1) &,& (0,1) \,.
\end{aligned}
\tag{52}
$$

The fluxes on M_1 and M_2 generate lower dimensional anti-brane like charges whereas the fluxes on M_3 generate lower dimensional brane like charges. The RR tadpole conditions for the $Z_2 \times Z_2$ orbifold without discrete torsion with $(\epsilon_1, \epsilon_2, \epsilon_3) = (-1, -1, 1)$ are given by

$$
\begin{aligned}
\sum_a M_a n_1^{(a)} n_2^{(a)} n_3^{(a)} = 16 &\quad,\quad \sum_a M_a n_1^{(a)} m_2^{(a)} m_3^{(a)} = 16 \,, \\
\sum_a M_a m_1^{(a)} n_2^{(a)} m_3^{(a)} = 16 &\quad,\quad \sum_a M_a m_1^{(a)} m_2^{(a)} n_3^{(a)} = -16 \,.
\end{aligned}
\tag{53}
$$

The massless spectrum in this class of models is determined by the intersection numbers

$$
I_{ab} = \prod_{ab} (n_i^{(a)} m_i^{(b)} - m_i^{(a)} n_i^{(b)}) \quad,\quad I_{ab} = \prod_{ab'} (n_i^{(a)} m_i^{(b)} + m_i^{(a)} n_i^{(b)}) \,,
$$
$$
I_{aO} = 8(m_i^{(a)} m_i^{(a)} m_i^{(a)} + m_i^{(a)} n_i^{(a)} n_i^{(a)} + n_i^{(a)} m_i^{(a)} n_i^{(a)} - n_i^{(a)} n_i^{(a)} m_i^{(a)}) \,.
\tag{54}
$$

The contribution of the four stacks of branes to the RR tadpole conditions with wrapping numbers (52) precisely satisfy (53) when $M_1 = M_2 = M_3 = M_4 = 4$. The gauge group of this model is $U(2)^3 \otimes SO(4)$. The model was chosen such that the chiral massless spectrum, determined by the intersection numbers, to contain only strings stretched between different stacks of branes. More precisely, by defining $M_i = 2p_i$, it is given by

$$
\phi_{1,\bar{a}b}^i \quad : \quad 16 \times (\bar{\mathbf{p}}_2, \mathbf{p}_3) \,, \quad \phi_{2,ac}^j \quad : \quad 16 \times (\mathbf{p}_1, \mathbf{p}_2) \,,
$$

$$\phi^k_{3,\bar{b}\bar{c}} \quad : \quad 16 \times (\bar{\mathbf{p}}_1, \bar{\mathbf{p}}_3) , \tag{55}$$

where the multiplicity of 16 in each sector comes from the intersection numbers of various branes, whereas all other charged states are non-chiral and will get a mass. There are mixed $U(1)_a \otimes U(1)_b^2$ and $U(1)_a \otimes SU(p_b)^2$ gauge anomalies in the model, easily computable from the massless spectrum.

$$C_{ab} = \frac{1}{4\pi^2} Tr(X_a X_b^2) = \frac{2^4}{4\pi^2} Tr(X_a T_b^2) = \frac{2^8}{4\pi^2} \begin{pmatrix} 0 & 1 & -1 \\ 1 & 0 & -1 \\ -1 & 1 & 0 \end{pmatrix} , \tag{56}$$

where X_a, $a = 1, 2, 3$ are the $U(1)_a$ gauge factors, whereas T_a are the nonabelian generators. They are taken care by axionic couplings of the type $\Theta_a F^a \wedge F^a$, where $\Theta_a = Im \ f_a$, where the gauge kinetic functions are given by

$$f_a = \prod_i n_i^{(a)} \ S - n_1^{(a)} m_2^{(a)} m_3^{(a)} \ T_1 - m_1^{(a)} n_2^{(a)} m_3^{(a)} \ T_2 - m_1^{(a)} m_2^{(a)} n_3^{(a)} \ T_3 \tag{57}$$

and where $Im \ S, Im \ T_i$ are the axion-dilaton and the three axions associated to the three internal tori. In our concrete example above, the gauge kinetic functions are explicitly $f_1 = S - 4T_1$, $f_2 = S - 4T_2$, $f_3 = S + 4T_3$. The mixed gauge anomalies are taken care by the nonlinear gauge transformations

$$\delta \ Im \ T_1 = \frac{16}{\pi^2} (-\alpha_2 + \alpha_3) \quad , \quad \delta \ Im \ T_2 = -\frac{16}{\pi^2} (\alpha_1 + \alpha_3)$$
$$\delta \ Im \ T_3 = -\frac{16}{\pi^2} (\alpha_1 + \alpha_2) , \tag{58}$$

where α_a are the gauge transformation parameters for the $U(1)_a$ factors. The Kähler potential contains the terms

$$- \ln \left[T_1 + \bar{T}_1 + \frac{16}{\pi^2} (V_2 - V_3) \right] - \ln \left[T_2 + \bar{T}_2 + \frac{16}{\pi^2} (V_1 + V_3) \right] - \ln \left(T_3 + \bar{T}_3 + \frac{16}{\pi^2} (V_1 + V_2) \right] \tag{59}$$

which generate FI terms in the effective field theory

$$\xi_1 = -\frac{16}{\pi^2} \left(\frac{1}{t_2} + \frac{1}{t_3} \right) , \quad \xi_2 = -\frac{16}{\pi^2} \left(\frac{1}{t_1} + \frac{1}{t_3} \right) , \quad \xi_3 = \frac{16}{\pi^2} \left(\frac{1}{t_1} - \frac{1}{t_2} \right) . \tag{60}$$

The Fayet-Iliopoulos terms can be more generally be written in terms of magnetic fluxes as

$$\xi_a \sim H_1^{(a)} + H_2^{(a)} + H_3^{(a)} - H_1^{(a)} H_2^{(a)} H_3^{(a)} \tag{61}$$

and in this model they satisfy the sum rule $\xi_1 - \xi_2 - \xi_3 = 0$. The D-terms on the U(1) factors of each $U(2) = U(1) \otimes SU(2)$ stack are given by

$$D_1 = \sum_{i,\bar{a},b} |\phi^i_{2,\bar{a}b}|^2 - \sum_{j,a,c} |\phi^j_{3,ac}|^2 + \xi_1 ,$$

$$D_2 = -\sum_{i,\bar{a},b} |\phi_{1,\bar{a}b}^i|^2 + \sum_{k,\bar{b},\bar{c}} |\phi_{2,b\bar{c}}^k|^2 + \xi_2 \ ,$$

$$D_3 = \sum_{j,a,c} |\phi_{1,ac}^j|^2 - \sum_{k,\bar{b},\bar{c}} |\phi_{3,\bar{b}\bar{c}}^k|^2 + \xi_3 \ , \tag{62}$$

and satisfy also the same rule

$$D_1 - D_2 - D_3 = \xi_1 - \xi_2 - \xi_3 = 0 \ . \tag{63}$$

We believe that the interpretation of the sum rule (63) is that the D-branes tend to recombine by condensing the by-fundamental fields (55) and to provide a supersymmetric vacuum $D_a = 0$. Notice however that the model has renormalizable superpotential terms

$$W = \lambda_{ijk} \, Tr \, (\phi_1^i \phi_2^j \phi_3^k) \ , \tag{64}$$

(where the trace in the gauge group space), which has a geometrical interpretation of Yukawa couplings connecting 3-fields forming a triangle in each compact torus, analogously to models of Yukawa couplings studied in [33]. The number of these Yukawas are related, as usual, to the number of D-flat directions. The Yukawa couplings are also field dependent and depend on complex structure moduli. In the presence of the Yukawa couplings (64), the tachyonic instabilities typically related to the brane recombination process can be removed, there is generically a geometrical obstruction and the D-brane recombination is not generically the most favorable process. The FI terms are Kähler moduli dependent and they can (perturbatively) vanish for particular points in the Kähler moduli space, which will always be dynamically preferred. In order to avoid this phenomenon, nonperturbative effects have to be invoked, for example gaugino condensation on the D9 branes, according to the discussion in the previous section. The perturbative nature of the superpotential terms (64), to be interpreted as meson masses in case ii) of the previous section, generate large D-term contributions . Cosmological constant cancellation for large D-terms in the TeV range gravitino mass ask for intermediate values of FI terms and/or very small $U(1)$ gauge couplings. So for large scale of supersymmetry breaking, cosmological constant is generically hard to cancel unless FI terms are much smaller than the Planck scale. Whereas this was not the case in the nonperturbative model of the previous section, the more general formula (61) suggests that it is possible that $\xi_a << H_i^{(a)}$ by tuning the magnetic fluxes, in the spirit of landscape models [9, 10, 22, 34, 35].

In order to make connection with the field theory model of the previous section, notice that if dynamics picks up an overall Kahler modulus $T_1 = T_2 = T_3 = T$, then $\xi_3 = 0$ and there are two remaining anomalous abelian factors $U(1)_1$, $U(1)_2$. In the following we discuss in some more detail a simplified model along these lines.

A. From intersecting brane models to nonperturbative moduli stabilisation

In order to stabilise all moduli, in section 4 we used nonpeturbative effects on an asymptotically-free gauge group. The explicit string example discussed previously has no asymptotically gauge factor, but we do not expect this result to be generic. In the following, we consider a model, similar to the explicit string example presented in the previous section but containing an asymptotically-free gauge factor. It is also further simplified in order to allow a simple analysis and suited to generate large D-terms at the minimum, comparable and in some regions of the parameter space dominant with respect to the F-terms.

The field content and gauge structure are summarized as follows. The model has a gauge group $SU(N) \otimes U(1)^2$, with the chiral superfield content

$$\phi_1^i : (N, 1, 0) \quad , \quad \phi_{2,\bar{j}} : (\bar{N}, 0, -1) \quad , \quad \phi_3 : (1, -1, +1) , \tag{65}$$

where $i, \bar{j} = 1 \cdots N_f$ and the notation for charges and representations are transparent in (65). The model is therefore similar to the explicit intersecting brane model of the previous section, but slightly adapted for our purposes. There is a magnetic flux pattern which does lead to the spectrum above, which by itself does not saturate the RR tadpole conditions. This can be cured by adding additional branes or by considering other orbifolds and/or additional antisymmetric field backgrounds. The $SU(N)$ plays the role of a hidden sector SYM gauge group with N_f flavors, which condenses in the IR. The composite objects

$$M_{\bar{j}}^i = \phi_1^{ia} \phi_{2,\bar{j}}^{\bar{a}}, \tag{66}$$

where a is an index in the fundamental of $SU(N)$, are the mesons used in constructing the effective action of the theory. For simplicity we consider in the following only the overall Kähler modulus T, whereas keeping all of them would ask for stabilisation a more complicated dynamics, for example several gaugino condensates. Consistently with the cancellation of the mixed gauge anomalies, the gauge kinetic function on the condensing gauge group is $f_{SU(N)} = S \pm N_f T$, where the $+$ $(-)$ signs correspond to a hidden sector with positive (negative) product of magnetic fluxes in two torii. Similar to the explicit intersecting brane model, T transforms under gauge transformations as

$$\delta T = \pm \frac{1}{4\pi^2} \left(\Lambda_2 - \Lambda_3 \right) . \tag{67}$$

This can also be directly checked by computing the mixed gauge anomalies

$$U(1)_2 \otimes SU(N)^2 \ : \ \frac{N_f}{4\pi^2} \quad , \ U(1)_3 \otimes SU(N)^2 \ : \ -\frac{N_f}{4\pi^2} , \tag{68}$$

which are precisely canceled by the nonlinear gauge transformation of the axion ImT (67). In order to write the Kähler potential, first of all we place ourselves on the $SU(N)$ flat direction $\langle \phi_1^{ia} \rangle = \langle \phi_{2,\bar{j}}^{\bar{a}} \rangle$. Similarly to the KKLT proposal, we could first integrate out the dilaton and the complex structure moduli. In doing this, for $N_f < N$, we find the effective superpotential and Kähler potential

$$W = W_0 + (N - N_f)A \left[\frac{e^{\mp 8\pi^2 N_f T}}{det M_{\bar{j}}^i} \right]^{\frac{1}{N - N_f}} + \lambda_{i\bar{j}} \phi_3 M_{\bar{j}}^i \ ,$$

$$K = -3 \ln \left[T + \bar{T} \pm \frac{V_3 - V_2}{4\pi^2} - 2 Tr(\bar{M}M)^{1/2} - \bar{\phi}_3 \phi_3 \right] \ , \tag{69}$$

where the constant W_0 depend on the details of the three-form fluxes, the Kähler potential was computed in the weakly coupled regime of the $SU(N)$ flat direction and where $A = \exp\{-8\pi^2 S_0/(N - N_f)\}$.

The reader will notice that this model reassembles closely the model worked out in Section 4. The mass term for the mesons in (69) is actually provided, in the intersecting brane realisation of the previous section, by the Yukawa coupling (64). In analogy with the explicit intersecting brane model, there is a constraint equation

$$D_2 - D_3 \ = \ \xi_2 - \xi_3 \ = \ 0 \ , \tag{70}$$

where

$$D_2 \ = \ Tr(\bar{M}M)^{1/2} - |\phi_3|^2 \pm \frac{3}{4\pi^2(T + \bar{T})} \ , \tag{71}$$

which signals the presence of a flat direction, allowing the presence of the perturbative superpotential providing the meson mass term. Indeed, since there are now two D-flatness conditions and two charged fields, the existence of the flat direction (70) is needed in order to write the meson mass term, the last term in the superpotential (69). The two signs in the expressions above correspond to the two cases $k > 0$ and $k < 0$ in Section 4, the second case realising the high scale supersymmetry breaking with large D-terms.

VII. SUMMARY AND OUTLOOK

In the present work, we have initiated a program to study in a general manner the implications of the large D-terms in a supergravity on soft supersymmetry breaking parameters. These terms can come from an anomalous $U(1)$ flavour model as has been noted in the past or tree level FI

terms in an intersecting D-brane model. We have shown that explicit models based on intersecting D-branes can be constructed giving rise to large D-terms. Irrespective of the source, we have studied the implications of these terms on the soft parameters. The mass squared terms are the most affected with contributions linear in D. However the charges of the matter fields under the anomalous $U(1)$ can crucially determine the actual impact. The B_μ and A terms also receive corrections though they are not significantly modified in terms of magnitude. As an application we have shown that split supersymmetry can be realised with specific choices of the superpotential and Kähler potential.

Particular examples with the large D contributions are string models of supersymmetry breaking, in particular in Type I string orientifolds. We have not addressed in detail the phenomenological signatures and constraints on the parameter space within this class of the models. Such a study could be taken in conjunction with a proper flavour model à la Froggatt-Nielsen. This could be then confronted with low energy data from accelerators, dark matter physics constraints and flavour physics.

The issue of moduli stabilisation has been receiving increasing attention in the recent years. Applications to the soft masses have also been recently addressed[37, 38]. Here we have revisited the issue of moduli stabilisation using non-perturbative gaugino condensation in type I orientifolds with internal magnetic fluxes which generate large D-terms. By a suitable choice of the fluxes which fix the sign of the FI term, we have shown that it is possible to stabilize moduli and generate large D-terms.

[1] R. Barbieri, S. Ferrara and C. A. Savoy, Phys. Lett. B **119**, 343 (1982); A. H. Chamseddine, R. Arnowitt and P. Nath, Phys. Rev. Lett. **49**, 970 (1982); L. J. Hall, J. Lykken and S. Weinberg, Phys. Rev. D **27**

(1983) 2359.

[2] H. P. Nilles, Phys. Rept. **110** (1984) 1.

[3] S. K. Soni and H. A. Weldon, Phys. Lett. B **126**, 215 (1983).

[4] V. S. Kaplunovsky and J. Louis, Phys. Lett. B **306**, 269 (1993) [arXiv:hep-th/9303040].

[5] A. Brignole, L. E. Ibanez and C. Munoz, Nucl. Phys. B **422**, 125 (1994) [Erratum-ibid. B **436**, 747 (1995)] [arXiv:hep-ph/9308271] ; A. Brignole, L. E. Ibanez and C. Munoz, arXiv:hep-ph/9707209.

[6] L. E. Ibanez, C. Munoz and S. Rigolin, Nucl. Phys. B **553**, 43 (1999) [arXiv:hep-ph/9812397].

[7] E. Dudas, S. Pokorski and C. A. Savoy, Phys. Lett. B **369**, 255 (1996) [arXiv:hep-ph/9509410; E. Dudas, C. Grojean, S. Pokorski and C. A. Savoy, Nucl. Phys. B **481**, 85 (1996) [arXiv:hep-ph/9606383]; Y. Kawamura and T. Kobayashi, Phys. Lett. B **375**, 141 (1996) [Erratum-ibid. B **388**, 867 (1996)] [arXiv:hep-ph/9601365] and Phys. Rev. D **56**, 3844 (1997) [arXiv:hep-ph/9608233].

[8] P. Binetruy and E. Dudas, Phys. Lett. B **389**, 503 (1996) [arXiv:hep-th/9607172]; G. R. Dvali and A. Pomarol, Phys. Rev. Lett. **77**, 3728 (1996) [arXiv:hep-ph/9607383].

[9] S. Weinberg, Phys. Rev. Lett. **59**, 2607 (1987); R. Bousso and J. Polchinski, JHEP **0006**, 006 (2000) [arXiv:hep-th/0004134].

[10] N. Arkani-Hamed and S. Dimopoulos, arXiv:hep-th/0405159; G. F. Giudice and A. Romanino, Nucl. Phys. B **699**, 65 (2004) [arXiv:hep-ph/0406088] .

[11] An incomplete list of references: L. Anchordoqui, H. Goldberg and C. Nunez, Phys. Rev. D **71**, 065014 (2005) [arXiv:hep-ph/0408284]; N. Arkani-Hamed, S. Dimopoulos, G. F. Giudice and A. Romanino, arXiv:hep-ph/0409232; D. A. Demir, arXiv:hep-ph/0410056; U. Sarkar, arXiv:hep-ph/0410104; R. Allahverdi, A. Jokinen and A. Mazumdar, Phys. Rev. D **71**, 043505 (2005) [arXiv:hep-ph/0410169]; E. J. Chun and S. C. Park, JHEP **0501**, 009 (2005) [arXiv:hep-ph/0410242]; B. Bajc and G. Senjanovic, Phys. Lett. B **610**, 80 (2005) [arXiv:hep-ph/0411193]; A. Arvanitaki and P. W. Graham, arXiv:hep-ph/0411376; A. Masiero, S. Profumo and P. Ullio, Nucl. Phys. B **712**, 86 (2005) [arXiv:hep-ph/0412058]; M. A. Diaz and P. F. Perez, J. Phys. G **31**, 563 (2005) [arXiv:hep-ph/0412066]; L. Senatore, Phys. Rev. D **71**, 103510 (2005) [arXiv:hep-ph/0412103]; A. Datta and X. Zhang, arXiv:hep-ph/0412255; M. Beccaria, F. M. Renard and C. Verzegnassi, Phys. Rev. D **71**, 093008 (2005) [arXiv:hep-ph/0412257]; P. C. Schuster, arXiv:hep-ph/0412263; S. P. Martin, K. Tobe and J. D. Wells, Phys. Rev. D **71**, 073014 (2005) [arXiv:hep-ph/0412424]; C. H. Chen and C. Q. Geng, arXiv:hep-ph/0501001; M. Drees, arXiv:hep-ph/0501106; S. Kasuya and F. Takahashi, arXiv:hep-ph/0501240; K. Huitu, J. Laamanen, P. Roy and S. Roy, arXiv:hep-ph/0502052; G. Marandella, C. Schappacher and A. Strumia, Nucl. Phys. B **715**, 173 (2005) [arXiv:hep-ph/0502095]; N. Haba and N. Okada, arXiv:hep-ph/0502213; C. H. Chen and C. Q. Geng, arXiv:hep-ph/0502246; B. Dutta and Y. Mimura, arXiv:hep-ph/0503052; D. Chang, W. F. Chang and W. Y. Keung, Phys. Rev. D **71**, 076006 (2005) [arXiv:hep-ph/0503055]; N. G. Deshpande and J. Jiang, Phys. Lett. B **615**, 111 (2005) [arXiv:hep-ph/0503116]; A. Ibarra, arXiv:hep-ph/0503160; B. Mukhopadhyaya and S. Sen-Gupta, arXiv:hep-ph/0503167; M. Toharia and J. D. Wells, arXiv:hep-ph/0503175; R. N. Mohapatra,

M. K. Parida and G. Rajasekaran, arXiv:hep-ph/0504236; J. G. Gonzalez, S. Reucroft and J. Swain, arXiv:hep-ph/0504260.

[12] E. Dudas and S. K. Vempati, arXiv:hep-ph/0506029.

[13] J. Wess and J. Bagger, "Supersymmetry and supergravity," Princeton, USA: Univ. Pr. (1992).

[14] S. Ferrara, C. Kounnas and F. Zwirner, Nucl. Phys. B **429**, 589 (1994) [Erratum-ibid. B **433**, 255 (1995)] [arXiv:hep-th/9405188].

[15] L. Randall and R. Sundrum, Nucl. Phys. B **557**, 79 (1999) [arXiv:hep-th/9810155]; G. F. Giudice, M. A. Luty, H. Murayama and R. Rattazzi, JHEP **9812**, 027 (1998) [arXiv:hep-ph/9810442].

[16] J. A. Bagger, T. Moroi and E. Poppitz, JHEP **0004**, 009 (2000) [arXiv:hep-th/9911029]; P. Binetruy, M. K. Gaillard and B. D. Nelson, Nucl. Phys. B **604**, 32 (2001) [arXiv:hep-ph/0011081].

[17] G. F. Giudice and A. Masiero, Phys. Lett. B **206**, 480 (1988).

[18] D. Z. Freedman, Phys. Rev. D **15**, 1173 (1977); P. Binetruy, G. Dvali, R. Kallosh and A. Van Proeyen, Class. Quant. Grav. **21**, 3137 (2004) [arXiv:hep-th/0402046].

[19] L. E. Ibanez and G. G. Ross, Phys. Lett. B **332**, 100 (1994) [arXiv:hep-ph/9403338]; P. Binetruy and P. Ramond, Phys. Lett. B **350**, 49 (1995) [arXiv:hep-ph/9412385]; E. Dudas, S. Pokorski and C. A. Savoy, Phys. Lett. B **356**, 45 (1995) [arXiv:hep-ph/9504292]. Y. Nir, Phys. Lett. B **354**, 107 (1995) [arXiv:hep-ph/9504312].

[20] Y. Nir and N. Seiberg, Phys. Lett. B **309**, 337 (1993) [arXiv:hep-ph/9304307]; M. Leurer, Y. Nir and N. Seiberg, Nucl. Phys. B **420**, 468 (1994) [arXiv:hep-ph/9310320].

[21] J. D. Wells, Phys. Rev. D **71**, 015013 (2005) [arXiv:hep-ph/0411041].

[22] I. Antoniadis and S. Dimopoulos, arXiv:hep-th/0411032; see also C. Kokorelis, arXiv:hep-th/0406258.

[23] P. J. Fox, A. E. Nelson and N. Weiner, JHEP **0208**, 035 (2002) [arXiv:hep-ph/0206096]; Z. Chacko, P. J. Fox and H. Murayama, Nucl. Phys. B **706**, 53 (2005) [arXiv:hep-ph/0406142].

[24] E. Cremmer, S. Ferrara, L. Girardello and A. Van Proeyen, Nucl. Phys. B **212** (1983) 413.

[25] G. R. Dvali and A. Pomarol, Phys. Lett. B **410**, 160 (1997) [arXiv:hep-ph/9706429].

[26] B. Kors and P. Nath, arXiv:hep-th/0411201; K. S. Babu, T. Enkhbat and B. Mukhopadhyaya, arXiv:hep-ph/0501079.

[27] For reviews, see for e.g, F. Quevedo, arXiv:hep-th/9603074 ; E. Dudas, Class. Quant. Grav. **17**, R41 (2000) [arXiv:hep-ph/0006190]; E. Kiritsis, Fortsch. Phys. **52**, 200 (2004) [arXiv:hep-th/0310001].

[28] M. Dine, N. Seiberg and E. Witten, Nucl. Phys. B **289**, 589 (1987); J. J. Atick, L. J. Dixon and A. Sen, Nucl. Phys. B **292**, 109 (1987).

[29] E. Witten, Nucl. Phys. B **202**, 253 (1982).

[30] T. R. Taylor, G. Veneziano and S. Yankielowicz, Nucl. Phys. B **218**, 493 (1983); I. Affleck, M. Dine and N. Seiberg, Nucl. Phys. B **241**, 493 (1984).

[31] C. Bachas, arXiv:hep-th/9503030; M. Berkooz, M. R. Douglas and R. G. Leigh, Nucl. Phys. B **480** (1996) 265 [arXiv:hep-th/9606139]; N. Ohta and P. K. Townsend, Phys. Lett. B **418** (1998) 77 [arXiv:hep-th/9710129];

[32] R. Blumenhagen, L. Goerlich, B. Kors and D. Lust, JHEP **0010** (2000) 006 [arXiv:hep-th/0007024]; C. Angelantonj, I. Antoniadis, E. Dudas and A. Sagnotti, Phys. Lett. B **489** (2000) 223 [arXiv:hep-th/0007090] ; G. Aldazabal, S. Franco, L. E. Ibanez, R. Rabadan and A. M. Uranga, JHEP **0102** (2001) 047 [arXiv:hep-ph/0011132]; M. Cvetic, G. Shiu and A. M. Uranga, Nucl. Phys. B **615** (2001) 3 [arXiv:hep-th/0107166] and Phys. Rev. Lett. **87**, 201801 (2001) [arXiv:hep-th/0107143].

[33] D. Cremades, L. E. Ibanez and F. Marchesano, JHEP **0307**, 038 (2003) [arXiv:hep-th/0302105]; M. Cvetic and I. Papadimitriou, Phys. Rev. D **68**, 046001 (2003) [Erratum-ibid. D **70**, 029903 (2004)] [arXiv:hep-th/0303083]; S. A. Abel and A. W. Owen, Nucl. Phys. B **663**, 197 (2003) [arXiv:hep-th/0303124].

[34] K. R. Dienes, E. Dudas and T. Gherghetta, arXiv:hep-th/0412185 ; N. Arkani-Hamed, S. Dimopoulos and S. Kachru, arXiv:hep-th/0501082.

[35] S. Kachru, R. Kallosh, A. Linde and S. P. Trivedi, Phys. Rev. D **68**, 046005 (2003) [arXiv:hep-th/0301240].

[36] C. P. Burgess, R. Kallosh and F. Quevedo, JHEP **0310**, 056 (2003) [arXiv:hep-th/0309187].

[37] K. Choi, A. Falkowski, H. P. Nilles, M. Olechowski and S. Pokorski, JHEP **0411**, 076 (2004) [arXiv:hep-th/0411066]; K. Choi, A. Falkowski, H. P. Nilles and M. Olechowski, Nucl. Phys. B **718**, 113 (2005) [arXiv:hep-th/0503216]; K. Choi, K. S. Jeong and K. i. Okumura, arXiv:hep-ph/0504037.

[38] J. P. Conlon, F. Quevedo and K. Suruliz, arXiv:hep-th/0505076.